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A REVISION OF: COSMIC MATTER ANTI MATTER ANNIHILATION AND THE γ-RAY BACKGROUND SPECTRUM

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COSMIC MATTER-ANTIMATTER ANNIHILATION AND THE γ-RAY BACKGROUND SPECTRUM

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ABSTRACT

We present here some initial results of a detailed calculation of the cosmological γ-ray spectrum from matter-antimatter annihilation in the universe. The similarity of the calculated spectrum with the present observations of the γ-ray background spectrum above 1 MeV suggests that such observations may be evidence of the existence of antimatter on a large scale in the universe. Quantitative comparison of the calculations with the existing observations indicates that the product of interacting matter and antimatter densities at present is $\tilde{n}_{p,0}\, n_{\bar{p},0} \lesssim 10^{-25}\ \mathrm{cm}^{-6}$.

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COSMIC MATTER-ANTIMATTER ANNIHILATION

AND THE γ-RAY BACKGROUND SPECTRUM

I. INTRODUCTION

The question of the existence of antimatter on a cosmological scale is one of the most basic problems of physics and cosmology. Recently it has taken on added interest because of the suggested possible role of antimatter in galaxy formation and the evolution of the universe (Alvén, 1965; Harrison, 1967; Omnès, 1969). It has long been recognized that the most promising way to search for evidence of antimatter on a cosmological scale is to attempt to observe γ-rays of cosmic origin which would be produced by the decay of neutral pions arising from matter-antimatter annihilation. In order to determine the annihilation-origin of such γ-rays, one must first calculate the γ-ray spectrum which would be produced by such annihilations so that the calculated spectrum may be compared with observational data on the γ-ray background spectrum for possible identification.

For the purpose of discussion, we define the terms "RG" and "CR" as follows. An RG (rest-gas) proton (or antiproton) is one possessing negligible kinetic energy with respect to the surrounding gas in an astronomical system. For example, nuclei of the galactic interstellar gas clouds from which stars are formed may be considered RG nuclei. A CR (cosmic-ray) proton (or antiproton) would be one possessing relativistic energy. We may thus envision four situations for discussion in which matter and antimatter might interact:

(1) RG matter + RG antimatter (RG-$\overline{\text{RG}}$)

(2) RG matter + CR antimatter (RG-$\overline{\text{CR}}$)

(3) CR matter + RG antimatter (CR-$\overline{\text{RG}}$)

(4) CR matter + CR antimatter (CR-$\overline{\text{CR}}$)

Although astrophysics has revealed many surprising phenomena, we would be hard put to imagine a system in which the number of RG nuclei did not greatly outnumber the number of CR nuclei. Indeed, the evolution of such a system from a more balanced one would violate the second law of thermodynamics. Therefore, if we let the number of RG nuclei in a typical system be N and the number of CR nuclei in the system be of the order ϵN, with $\epsilon << 1$, then typical interaction rates would have the properties

$$\frac{R(\text{RG}-\overline{\text{CR}})}{R(\text{RG}-\overline{\text{RG}})} \sim \frac{R(\text{CR}-\overline{\text{RG}})}{R(\text{RG}-\overline{\text{RG}})} \sim \epsilon \tag{1}$$

$$\frac{R(\text{CR}-\overline{\text{CR}})}{R(\text{CR}-\overline{\text{RG}})} \sim \frac{R(\text{CR}-\overline{\text{CR}})}{R(\text{RG}-\overline{\text{CR}})} \sim \epsilon \tag{2}$$

and

$$\frac{R(\text{CR}-\overline{\text{CR}})}{R(\text{RG}-\overline{\text{RG}})} \sim \epsilon^2 \tag{3}$$

We will thus assume that the likelihood of CR-$\overline{\text{CR}}$ interactions is negligible compared to the likelihood of RG-$\overline{\text{CR}}$ or CR-$\overline{\text{RG}}$ interactions that will produce γ-rays of similar characteristics, and we may limit ourselves to the discussion of RG-$\overline{\text{RG}}$ and RG-$\overline{\text{CR}}$ (or CR-$\overline{\text{RG}}$) interactions.

II. THE γ-RAY SOURCE SPECTRUM

The γ-ray source spectrum from p-$\bar{\text{p}}$ interactions of the RG-$\overline{\text{CR}}$ type is given by

$$q_{RG\text{-}\overline{CR}}(E_\gamma, \vec{r}) = 4\pi n_p(\vec{r}) \int dE_{\bar{p}}\, I(E_{\bar{p}}, \vec{r}) \sum_s \int dE_s\, \sigma_s(E_s|E_{\bar{p}})$$

$$\times \sum_d \zeta_{\gamma d} R_{\gamma d} f_{ds}(E_\gamma|E_s) \tag{4}$$

For interactions of the CR-$\overline{\text{RG}}$ type, we use the corresponding expression

$$q_{CR\text{-}\overline{RG}}(E_\gamma, \vec{r}) = 4\pi n_{\bar{p}}(\vec{r}) \int dE_p\, I(E_p, \vec{r}) \sum_s \int dE_s\, \sigma_s(E_s|E_p)$$

$$\times \sum_d \zeta_{\gamma d} R_{\gamma d} f_{ds}(E_\gamma|E_s) \tag{5}$$

In equations (4) and (5), the quantity $n(\vec{r})$ is the number of target nucleons in the medium per cubic centimeter as a function of position; $I(E, \vec{r})$ is the differential cosmic-ray particle flux in cm^{-2}-s^{-1}-sr^{-1}-GeV^{-1}. The subscript p stands for proton, $\bar{p}$ for antiproton, s for secondary particle produced in the collision, and d for decay mode. The production function $\sigma_s(E_s|E_p)$ represents the cross section for production of secondary particles of type s and energy E_s in a collision of primary energy E_p, $\zeta_{\gamma d}$ represents the number of γ-rays produced in the decay mode d, $R_{\gamma d}$ is the branching ratio for the decay mode d (the probability that a secondary particle s will decay via mode d) and $f_{ds}(E_\gamma|E_s)$ is the normalized distribution function representing the probability that a secondary particle with energy E_s will decay to produce a γ-ray of energy E_γ. For the purposes of discussion, we will restrict ourselves to p-$\bar{p}$ interactions.

The situation for matter-antimatter annihilations occurring near rest is somewhat different. In this case, we will find it more convenient to speak of particle densities instead of fluxes or intensities. We thus make the transformation

$$\int dE_k\, I_k(E_k, \vec{r}) \rightarrow \frac{n_k(\vec{r})}{4\pi} \int dv\; v f(v) \;=\; \frac{n_k(r) <v>}{4\pi} \tag{6}$$

where v is the relative velocity between the annihilating nucleon and antinucleon, which we can consider to be a thermal velocity, and f (v) is a normalized distribution function representing the distribution of relative velocities between the interacting nucleons.

The γ-ray source spectrum from p-$\overline{\text{p}}$ interactions at or near rest is given by

$$q_{RG-\overline{RG}}(E_\gamma, \vec{r}) = n_p(\vec{r}) n_{\overline{p}}(\vec{r}) \int dv\, f(v)\, v \sum_s \int dE_s\, \sigma_s(E_s|v) \times \sum_d \zeta_{\gamma d} R_{\gamma d} f_{ds}(E_\gamma|E_s) \tag{7}$$

We therefore define the emission measure B as

$$B = \int dl\, n_p(\vec{r}) n_{\overline{p}}(\vec{r}) \tag{8}$$

It follows from equation (7) that

$$I_{RG-RG}(E_\gamma) = \frac{1}{4\pi} \int dl\, q(\vec{r}) \propto B \tag{9}$$

where l is the distance along the line of sight.

III. CROSS SECTIONS FOR PROTON-ANTIPROTON ANNIHILATIONS NEAR REST

We now state more precisely what we mean by a "rest gas." We define a rest gas to be a gas of particles such that no particle in the gas has an energy greater than 286 MeV. This rather liberal definition of rest is sufficient to

insure that the only secondary particles produced in RG-$\overline{\text{RG}}$ interactions that yield γ-rays are secondary mesons produced by nucleon-antinucleon annihilation. Our restriction leaves out interactions of the type

$$p + \bar{p} \to p + \bar{p} + \pi^0 \tag{10}$$

since the threshold for reactions of this type is 286 MeV.

Since the threshold for nonannihilation, inelastic, p-$\bar{p}$ interactions that produce particles other than π^0-mesons is even greater, it follows that in the RG-$\overline{\text{RG}}$ case, only annihilations; i.e., reactions of the form

$$\begin{array}{l} p + \bar{p} \to \text{bosons} \\ \qquad\qquad \hookrightarrow \text{neutral pions} \\ \qquad\qquad\qquad\qquad \hookrightarrow \gamma\text{-rays} \end{array} \tag{11}$$

need be considered.

Table 1 lists the experimental cross sections for free p-$\bar{p}$ annihilation as a function of incident antiproton kinetic energy at accelerator energies. Also listed are β_{cms} and the product of the annihilation cross section (σ_A) and cms velocity (β_{cms}). It can be seen from table 1 that the product of cms velocity and annihilation cross section is constant over a very large energy range. For nonrelativistic energies,

$$\beta_{cms} \simeq \frac{1}{2}\beta \qquad (\beta << 1) \tag{12}$$

β being the relative velocity of the particles. We find experimentally that

$$\sigma_A \simeq \frac{4.8 \times 10^{-26} \ \text{cm}^2}{\beta} \simeq \frac{1.4 \times 10^{-15} \ \text{cm}^3\text{-s}^{-1}}{v} \tag{13}$$

At all energies above 25 MeV, we find that table 1 agrees with the theoretical cross section

$$\sigma_A = \frac{\pi r_p^2}{\beta_{cms}} \tag{14}$$

taking

$$r_p = 0.87 \times 10^{-13} \text{ cm} \tag{15}$$

as the proton (or antiproton) radius and using the model of Koba and Takeda (1958) in which the nucleon acts as a black absorbing sphere of radius r_p.

At lower energies ($\lesssim$ 10 MeV, which correspond to annihilations of matter and antimatter gases at a temperature $\lesssim 10^{11}$ K), coulomb forces play a dominant role in the matter-antimatter interaction process. At energies $\lesssim$ 100 eV, bound systems of protons and antiprotons can be formed as an intermediate stage before annihilation. The situation is similar to the case of positronium formation (Stecker 1969). As in that case, we find that under astrophysical conditions, the lifetime of the p-$\bar{p}$ system against annihilation is much shorter than its lifetime against breakup, so that the cross section for bound-state formation becomes the effective annihilation cross section. Until recently the effect of coulomb interactions on matter-antimatter annihilation was largely neglected and theoretical cross sections for bound-state formation processes had not been calculated. However, recently Morgan and Hughes (1970) have calculated the cross sections for the radiative capture processes

$$e^- + e^+ \rightarrow \Pi_e + \gamma \tag{16}$$

and

$$p + \bar{p} \rightarrow \Pi_p + \gamma \tag{17}$$

and the atomic processes

$$\bar{H} + H \rightarrow \Pi_p + (\Pi_e \text{ or } e^- + e^+) \tag{18}$$

$$H_2 + \bar{H} \rightarrow \Pi_p + (\Pi_e \text{ or } e^- + e^+) + (H \text{ or } p + e^-) \tag{19}$$

$$H_2 + \bar{H}_2 \rightarrow 2\Pi_p + (\text{various combinations of } 2e^-,\ 2e^+) \tag{20}$$

$$H_2 + \bar{H}_2 \rightarrow \Pi_p + (\Pi_e \text{ or } e^+ + e^-) + (\text{various combinations of } p,\ \bar{p},\ e^+, \text{ and } 1e^-) \tag{21}$$

$$p + \bar{H} \rightarrow \Pi_p + e^+ \tag{22}$$

$$p + \bar{H}_2 \rightarrow \Pi_p + (\Pi_e \text{ or } e^+ + e^-) + (\bar{H} \text{ or } \bar{p} + e^+) \tag{23}$$

$$e^- + \bar{H} \rightarrow \Pi_e + \bar{p} \tag{24}$$

and

$$e^+ + H \rightarrow \Pi_e + p \tag{25}$$

where we have used the symbols Π_p and Π_e to denote the bound states of the p-$\bar{p}$ and e^--e^+ systems, respectively. Reactions (24) and (25) have already been discussed with regard to the annihilation of secondary cosmic-ray positrons in a tenuous neutral hydrogen gas (Stecker 1969). The cross sections for these reactions were calculated by Cheshire (1964). With regard to p-$\bar{p}$ annihilation in intergalactic space, three processes in particular can be expected to be of the greatest importance:

(1) Direct annihilations, such as (11), at energies greater than 10 MeV, where the cross section can be well represented by equations (14) and (15).

(2) Direct annihilations at energies below 10 MeV, where the mutual coulomb attraction of the proton and antiproton distort the wave functions

of the two particles with a resultant increase in the direct-annihilation cross section. When the cross section given by equation (14) is corrected for this effect, the cross-section formula becomes modified to

$$\sigma_A = \frac{2\pi\alpha/\beta}{1 - \exp(-2\pi\alpha/\beta)} \frac{\pi r_p^2}{\beta_{cms}} \tag{26}$$

where α is the fine structure constant. For p-$\overline{\text{p}}$ interactions at energies much less than 4 MeV, where $\beta << 2\pi\alpha$, equation (26) reduces to

$$\sigma_A \simeq \frac{4\pi^2 \alpha r_p^2}{\beta^2}, \qquad \beta << 2\pi\alpha \tag{27}$$

Thus, at an energy of about 4 MeV, there is a transition in the velocity dependence of the annihilation cross section from $\sigma_A \sim \beta^{-1}$ to $\sigma_A \sim \beta^{-2}$. The same is true for electron-positron interactions at energies of the order of 2 keV where $\beta \approx 2\pi\alpha$.

(3) In interactions between neutral atomic gases of hydrogen and antihydrogen at thermal energies, the H-$\overline{\text{H}}$ rearrangement collision, reaction (18), becomes the dominant mode of p-$\overline{\text{p}}$ annihilation. Morgan and Hughes (1970) have calculated the annihilation cross section for this reaction and found that for energies between 10^{-3} and 1 eV (corresponding to thermal velocities at temperatures between 10 K and 10^4 K)

$$\sigma_{H\text{-}\overline{H},A} \simeq (0.31\, a_0)^2\, \beta^{-0.64} \tag{28}$$

where a_0 is the Bohr radius of the hydrogen atom.

Radiative capture reactions of the form of equation (17) can be shown to be insignificant. The cross section for annihilations of this type (Morgan and Hughes, 1970) is

$$\sigma_{A,r} = \frac{128}{3\sqrt{3}} \alpha \left(\frac{m_e}{m_p}\right)^2 \left[\ln\left(\frac{\alpha}{\beta}\right) + 0.20 + 0.25\left(\frac{\alpha}{\beta}\right)^{-2/3}\right] \frac{\sigma_0}{\beta^2} \tag{29}$$

Equation (29) is accurate to within at least 1 percent when $\beta \leq \alpha/7$ and is exact in the limit as $\beta \to 0$. The radiative capture cross section for the electron-positron system, viz, $e^+ + e^- \to \Pi_e + \gamma$, is given by the same formula without the quantity $(m_e/m_p)^2$. For p-$\bar{p}$ annihilations at all energies of interest $\sigma_{A,r} << \sigma_A$; however, in the case of electron-positron annihilation (in an ionized medium) $\sigma_{A,r} > \sigma_A$ for interactions at energies less than 20 eV.

In figure 1 we give the cross section versus temperature, energy, and thermal velocity for the reactions expected to be of importance for cosmological problems (intergalactic p-$\bar{p}$ annihilations). Figure 2 shows the values for total-annihilation cross section times velocity (proportional to the annihilation rate) based on the calculations of Morgan and Hughes, which should be valid for astrophysical and cosmological problems at the temperatures and kinetic energies indicated.*

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*In interactions between partially ionized atomic gases of hydrogen and antihydrogen, reaction (22) and its charge conjugate analogue should also be considered. Morgan and Hughes find the cross section for this reaction to be

$$\sigma_{p+\bar{H},\bar{p}+H} \simeq (0.057\, a_0)^2\, \beta^{-1}$$

for temperatures between 0.1 and 3000K. When partial ionization as a function of temperature is taken into account using the Saha equation, it is found that the total effective cross section is well approximated by (28) for temperatures $\lesssim$ 5000K and by (26) or (27) for greater temperatures.

We may now make use of equations (17), (18), (27), and (28) in order to reduce equation (7) to a simpler form. To do this, let us assume that matter-antimatter annihilations are taking place at thermal energies in three temperature regions defined as

$$
\begin{aligned}
&\text{Region I:} && 10^{11}\ \text{K} \lesssim T \lesssim 10^{13}\ \text{K} \\
&\text{Region II:} && 10^{4}\ \text{K} \lesssim T \lesssim 10^{11}\ \text{K} \\
&\text{Region III:} && 10\ \text{K} \lesssim T \lesssim 10^{4}\ \text{K}
\end{aligned}
\tag{30}
$$

It follows from the equations cited above that in each of these regions the annihilation cross section varies as some inverse power of the velocity. We can write an expression for this cross section for each of the temperature regions defined above thus

$$
\begin{aligned}
&\text{Region I:} && \sigma(E_s|v) \simeq \sigma_I \left(\frac{v}{c}\right)^{-1} \\
&\text{Region II:} && \sigma(E_s|v) \simeq \sigma_{II} \left(\frac{v}{c}\right)^{-2} \\
&\text{Region III:} && \sigma(E_s|v) \simeq \sigma_{III} \left(\frac{v}{c}\right)^{-0.64}
\end{aligned}
\tag{31}
$$

where

$$
\begin{aligned}
\sigma_I &= 4.8 \times 10^{-26}\ \text{cm}^2 \\
\sigma_{II} &= 2.2 \times 10^{-27}\ \text{cm}^2
\end{aligned}
\tag{32}
$$

and

$$\sigma_{III} = 2.6 \times 10^{-18}\ \text{cm}^2$$

We will write for all three regions

$$\sigma(E_s|v) = \sigma_i \left(\frac{v}{c}\right)^{-\delta_i} f(E_s) \tag{33}$$

where δ_i is the exponent of the power-law velocity dependence of $\sigma(E_s|v)$ as given in equation (31) and where, as in previous discussions

$$\int_0^\infty dE_s\, f(E_s) \equiv 1 \tag{34}$$

If we assume that the velocity distribution of the interacting nucleons and antinucleons is thermal and can be represented by the normalized maxwellian distribution

$$f_M(v) = \sqrt{\frac{2}{\pi}} \left(\frac{m}{kT}\right)^{3/2} v^2 \exp\left(-\frac{mv^2}{2kT}\right) \tag{35}$$

at some given temperature T; and in addition, if we define a normalized energy spectrum of the γ-radiation in the rest frame of the gas,

$$\zeta_\gamma f(E_{\gamma 0}) \equiv \sum_s \int dE_s\, f(E_s) \sum_d \zeta_{\gamma d} R_{\gamma d} f_{ds}(E_{\gamma 0}|E_s) \tag{36}$$

and smear it out by the distribution function,

$$D(E_\gamma|E_{\gamma 0}) \equiv \left(\frac{mc^2}{2\pi kT}\right)^{1/2} \exp\left[-\frac{mc^2}{2kT}\left(\frac{E_\gamma - E_{\gamma 0}}{E_{\gamma 0}}\right)^2\right] \tag{37}$$

to take account of the thermal Doppler broadening of the γ-ray spectrum, we may write equation (7) in the form

$$q_{RG\text{-}\overline{RG}}(E_\gamma, \vec{r}) = n_p(\vec{r})\, n_{\bar{p}}(\vec{r})\, \sigma_i\, c^{\delta_i}\, \zeta_\gamma \left(\frac{2}{\pi}\right)^{1/2} \left(\frac{m}{kT}\right)^{3/2}$$

$$\times \int dv\, v^{(3-\delta_i)} \exp\left(-\frac{mv^2}{2kT}\right) \int dE_{\gamma 0}\, D(E_\gamma | E_{\gamma 0})\, f(E_{\gamma 0}) \qquad (38)$$

Equation (38) may be further reduced by elimination of the first integral to read

$$q_{RG\text{-}\overline{RG}}(E_\gamma, r) = n_p(\vec{r})\, n_p(\vec{r})\, \sigma_i\, c^{\delta_i}\, \zeta_\gamma \frac{2}{\pi^{1/2}}\, \Gamma\left(2 - \frac{\delta_i}{2}\right)$$

$$\times \left(\frac{2kT}{m}\right)^{1/2(1-\delta_i)} \int dE_{\gamma 0}\, D(E_\gamma | E_{\gamma 0})\, f(E_{\gamma 0}) \qquad (39)$$

where $\Gamma(x)$ is the gamma function.

In region I, $\delta_I = 1$, and therefore, from equation (7),

$$\int dv\, f(v)\, v\sigma_I c v^{-1} = \sigma_I c \int dv\, f(v)\, v \cdot v^{-1} = \sigma_I c \int dv\, f(v)$$

$$= \sigma_I c \qquad (40)$$

since $\int dv\, f(v) \equiv 1$ by definition.

We therefore find that in region I,

$$q_{RG\text{-}\overline{RG}}(E_\gamma, \vec{r}) = n_p(\vec{r})\, n_{\bar{p}}(\vec{r})\, \sigma_I c \zeta_\gamma \int dE_{\gamma 0}\, D(E_\gamma | E_{\gamma 0})\, f(E_{\gamma 0})$$

$$\text{for } 10^{11}\ \text{K} \lesssim T \lesssim 10^{13}\ \text{K} \qquad (41)$$

In temperature regions II and III, we cannot make the simplification given by equation (40); however, we can make a different simplification. In these

temperature regions, $\beta << 1$, and Doppler broadening due to thermal effects provides only a negligible distortion on the γ-ray spectrum. Thus, in these temperature regions, we may write

$$D(E_\gamma|E_{\gamma 0}) \simeq \delta(E_\gamma - E_{\gamma 0}) \qquad (\beta << 1) \tag{42}$$

and therefore

$$\int dE_{\gamma 0} D(E_\gamma|E_{\gamma 0})\, f(E_{\gamma 0}) \doteqdot f(E_\gamma) \tag{43}$$

In region II, $\delta_{II} = 2$, and by reducing equation (39) and making use of the approximation (43), we obtain

$$q_{RG-\overline{RG}}(E_\gamma, \vec{r}) = n_p(\vec{r})\, n_{\bar{p}}(\vec{r})\, \sigma_{II}\, c\, \zeta_\gamma \left(\frac{2mc^2}{\pi kT}\right)^{1/2} f(E_\gamma)$$

$$\text{for } 10^4 \text{ K} \lesssim T \lesssim 10^{11} \text{ K} \tag{44}$$

In region III, $\delta_{III} = 0.64$, and again reducing equation (39) and using the approximation (43), we find

$$q_{RG-\overline{RG}}(E_\gamma, \vec{r}) = n_p(\vec{r})\, n_{\bar{p}}(\vec{r})\, \sigma_{III}\, c\, \zeta_\gamma \left[\frac{2}{\pi^{1/2}}\Gamma(1.68)\left(\frac{2kT}{mc^2}\right)^{0.18}\right] f(E_\gamma)$$

$$\text{for } 10 \text{ K} \lesssim T \lesssim 10^4 \text{ K} \tag{45}$$

IV. MESON AND γ-RAY PRODUCTION IN p-$\bar{p}$ ANNIHILATIONS

Both the experimental data and the simple statistical model of Matsuda (1966) indicate that in p-$\bar{p}$ annihilations at rest, ρ-meson production is an order of magnitude less important than pion production and that production of other

mesons is at least two orders of magnitude less frequent than pion production, (see table 2) Table 3 shows the decay schemes of these mesons that lead to final-state γ-rays and other relevant data. Table 4 shows some recent data on meson production indicating that about 20 percent of the γ-rays produced arise through nonpionic meson production. The largest nonpion contribution to the γ-ray spectrum is due to the ρ-meson decay schemes

$$\rho^{\pm} \rightarrow \pi^{\pm} + \underbrace{\pi^{0}}_{\rightarrow \gamma + \gamma} \tag{46}$$

The ρ meson is an isospin triplet (T = 1) constructed from two pions, each having T = 1. Evaluation of the Clebsch-Gordon coefficients for this construction yields

$$\begin{aligned} |\rho^{+}\rangle &= \frac{1}{\sqrt{2}}\left(|\pi^{+}\rangle|\pi^{0}\rangle - |\pi^{0}\rangle|\pi^{+}\rangle\right) \\ |\rho^{0}\rangle &= \frac{1}{\sqrt{2}}\left(|\pi^{+}\rangle|\pi^{-}\rangle - |\pi^{-}\rangle|\pi^{+}\rangle\right) \\ |\rho^{-}\rangle &= \frac{1}{\sqrt{2}}\left(|\pi^{0}\rangle|\pi^{-}\rangle - |\pi^{-}\rangle|\pi^{0}\rangle\right) . \end{aligned} \tag{47}$$

Since the ρ^0 construction does not contain any $|\pi^0\rangle|\pi^0\rangle$ terms,

$$|\langle\pi^0\pi^0|\rho^0\rangle|^2 = 0 \tag{48}$$

and, therefore,

$$\rho^0 \not\rightarrow \pi^0 + \pi^0 \tag{49}$$

Gamma-rays from $\rho^{\pm}$ decay (reaction (46)) possess an average energy of 210 MeV, not much different from the 190-MeV average energy given to γ-rays from the directly produced pions that are an order of magnitude more

frequent. Because other mesons are produced even less frequently than the ρ mesons, we can conclude that mesons other than pions have a negligible effect on the total γ-ray spectrum from p-$\overline{\text{p}}$ annihilation.

If we assume that the large majority of γ-rays from RG-$\overline{\text{RG}}$ interactions arise through π^0 decay, equation (36) reduces to

$$f(E_\gamma) = \int_{E\gamma + (m_\pi^2/4E\gamma)}^{\infty} dE_\pi \frac{f_A(E_\pi)}{(E_\pi^2 - m_\pi^2)^{1/2}} \tag{50}$$

The process

$$p + \overline{p} \rightarrow \pi^0 \tag{51}$$

is, of course, forbidden by conservation of momentum. The process

$$p + \overline{p} \rightarrow \pi^0 + \pi^0 \tag{52}$$

is also forbidden since p-$\overline{\text{p}}$ annihilations at rest occur predominantly from the S states of the p-$\overline{\text{p}}$ system.

The selection rule that forbids reaction (52) follows from conservation of G conjugation parity.

The process of G conjugation is an extension of charge (C) conjugation, which holds for neutral as well as charged particles. It is defined as

$$G = Ce^{i\pi T_2} \tag{53}$$

where C is the charge-conjugation operator and T_2 is the second component of the isospin vector. From (53) we can show the commutation relation

$$[T, G] = 0 \tag{54}$$

Thus, we may describe particle states as simultaneous eigenstates of both G and T.

It can be shown (Sakurai, 1964) that systems having baryon number 0 are in an eigenstate of G. The p-$\overline{p}$ system is just such a system. For this system

$$G = (-1)^{L+S+T} \tag{55}$$

where L, S, and T are the orbital, spin, and isospin quantum numbers of the state, respectively.

In a state consisting of a single pion, L = 0, S = 0, T = 1, and, therefore, G = - 1. In a final state consisting of ζ_π pions, G is given by

$$G = (-1)^{\zeta_\pi} \tag{56}$$

The selection rules (55) and (56) indicate that for an S-state annihilation (L = 0), the final state consisting of two neutral pions is strictly forbidden (Lee and Yang, 1956).

Therefore, the extremum γ-ray energies that we would expect from this decay would result from interactions of the type

$$p + \overline{p} \rightarrow \pi^+ + \pi^- + \pi^0 \tag{57}$$

resulting in pions with the given maximum energy

$$E_{\pi^0, max} = \frac{1}{2(2m_p)} \left[(2m_p)^2 + m_{\pi^0}^2 - (2m_{\pi\pm})^2 \right] \tag{58}$$

$$= m_p - \frac{3}{4} \frac{m_\pi^2}{m_p} \simeq 923 \text{ MeV}$$

Thus, the annihilation γ-rays are limited to the energy region

$$\frac{1}{2}\left[E_{\pi^0,\,min} - (E^2_{\pi^0,\,max} - m_\pi^2)^{1/2}\right] \leq E_\gamma \leq \frac{1}{2}\left[E^0_{\pi,\,max} + (E^2_{\pi^0,\,max} - m_\pi^2)^{1/2}\right] \tag{59}$$

or 5 MeV $\leq E_\gamma \leq$ 919 MeV.

The process

$$p + \overline{p} \rightarrow \gamma + \gamma \tag{60}$$

may also occur; but this process, involving the electromagnetic emission (em) of two photons, is of the order

$$\sigma_{em} = \alpha^2 \sigma_{A,\,strong} \simeq \frac{2.6 \times 10^{-30}\ \mathrm{cm}^2}{\beta} \tag{61}$$

and therefore this process may be neglected as a significant contribution to γ-ray production.

V. THE GAMMA-RAY SPECTRUM FROM PROTON-ANTIPROTON ANNIHILATIONS AT REST

As we showed in the previous section, we can neglect the contribution from the decay of mesons other than neutral pions in calculations of the γ-ray spectrum from p-$\overline{p}$ annihilations at rest. The normalized γ-ray spectrum was calculated numerically from the relation (50) with the normalized distribution function $f_A(E_\pi)$ taken from the calculations of Maksimenko (1958), based on the statistical theory of multiple particle production.

The resultant spectrum, up to 750 MeV, is shown in figure 3. Frye and Smith (1966) have recently calculated the γ-ray spectrum from p-$\overline{p}$ annihilation up to 500 MeV, based on recent measurements of charged pions from p-$\overline{p}$ annihilation by the Columbia University group. The excellent agreement between

the results of figure 3 and the calculations of Frye and Smith not only serves as a mutual check on the calculations but also supports our previous conclusion that mesons other than pions have a negligible effect on the total γ-ray spectrum from p-$\bar{p}$ annihilation at rest.

VI. COSMOLOGICAL REDSHIFTED SPECTRA

Our calculations are in accord with the suggestions proposed by the models of Harrison[2] and Omnes[3]. These models assume that the universe consists of equal amounts of matter and antimatter which were separated into distinct regions at the earliest stages in the big-bang model of the universe. In the appendix it is shown that in the energy region of observational interest discussed here ($E_\gamma \gtrsim 1$ MeV), where the redshift of origin of the γ-ray background is $\lesssim 100$, we can assume that annihilations take place on the boundaries of colliding regions of matter and antimatter so that the total annihilation rate is proportional to $(1 + z)^6$. We use the commonly defined cosmological parameters H_0 and Ω where H_0 is the Hubble constant and Ω is the ratio of the average atomic matter density in the universe to the critical density $n_c \simeq 10^{-5}$ cm^{-3} needed to gravitationally close the universe. We also define the quantity ξ which denotes the mean fraction of the total atomic density interacting at the boundaries between the regions of matter and antimatter. The atomic densities n_p and $n_{\bar{p}}$ are each proportional to $(1 + z)^3$ and it is assumed that on the average $n_p = n_{\bar{p}} = \Omega n_c$. Thus we can define a mean density of interacting matter (or antimatter) $\tilde{n}^{int}$ given by $\tilde{n}^{int} = \xi \Omega n_c$.

For cosmological γ-ray production, the quantity B defined by equation (8) then becomes

$$B \to \xi^2 \int dz\, n_p(z)\, n_{\bar{p}}(z) \frac{dl}{dz} \tag{62}$$

z is defined as the redshift suffered by a photon originating at a distance l.

The derivative

$$\frac{dl}{dz} = \frac{c}{H_0(1+z)^2(1+\Omega z)^{1/2}}$$

where

$$n_p(z) = n_{p,0}(1+z)^3$$

and (63)

$$n_{\bar{p}}(z) = n_{\bar{p},0}(1+z)^3$$

(see, e.g., McVittie 1965).

Defining

$$B_0 \equiv \xi^2\left(\frac{c}{H_0}\right)n_{p,0}\,n_{\bar{p},0} \tag{64}$$

we find by taking account of the effects of redshift, expansion, and time dilatation in the expanding universe

$$I_A(E_\gamma) = \frac{B_0\sigma_i c^{\delta_i}}{2\pi^{3/2}}\left(\frac{2k}{m}\right)^{\left(\frac{1-\delta_i}{2}\right)}\Gamma\left(2-\frac{\delta_i}{2}\right)$$

$$\times\int dz\,[T(z)]^{\left(\frac{1-\delta_i}{2}\right)}\frac{(1+z)^2}{(1+\Omega z)^{1/2}}G_A[(1+z)E_\gamma] \tag{65}$$

where i = II, III; $\sigma_{II} = 2$; $\delta_{III} = 0.64$; $\sigma_{II} = 1.1 \times 10^{-27}$ cm^2; $\sigma_{III} = 2.6 \times 10^{-18}$ cm^2; $G_A(E_\gamma) \equiv \zeta_\gamma f_A(E_\gamma)$ where $f_A(E_\gamma)$ is shown in figure 3; and T_m is the matter temperature of the universe, which is a function of z. The matter temperature

$$T(z) \simeq \begin{cases} 2.7\mathrm{K}\,(1+z) & \text{for } z \gtrsim 150 \\ 1.8\times 10^{-2}\mathrm{K}\,(1+z)^2 & \text{for } z \lesssim 150 \end{cases} \qquad (66)$$

(Zel'dovich, et al., 1969).

It follows that equations (65) and (66) together with the bounded form of the function $G_A(E_\gamma)$ yield source spectra of the form $E^{-\Gamma}$ as shown in table 5 when absorption effects can be neglected.

At high redshifts, when pair production and Compton scattering become important, it becomes necessary to solve a cosmological-photon-transport (CPT) equation in order to determine the γ-ray spectrum.

VII. THE COSMOLOGICAL-PHOTON-TRANSPORT EQUATION

The cosmological-photon-transport (CPT) equation for x-rays and γ-rays has been discussed at great length by Arons (1971a, b). Our formulation differs from his basically only in the definition of the γ-ray intensity, I. Our intensity is defined as a photon intensity in units of photons per cm^2-s-sr-MeV (or in energy units of 0.511 MeV = m_ec^2). The intensity defined by Arons is in terms of an energy flux rather than a photon flux. The energy flux decreases more sharply with redshift than the photon flux, since the energy of each photon suffers an adiabatic expansion loss as the universe expands. Thus, the form of the CPT equation used here is slightly different than that used by Arons.

Defining $y \equiv 1 + z$, and $\epsilon = E_\gamma/m_ec^2$, we can write the CPT equation as follows

$$\frac{\partial}{\partial t}\left(y^{-3}\, I(\epsilon,y)\right) + \frac{\partial}{\partial \epsilon}\left(\frac{d\epsilon}{dt}\, y^{-3}\, I(\epsilon,y)\right) = y^{-3}\, Q_A(\epsilon,y)$$

$$- n(y)\, \sigma_a(\epsilon)\, c y^{-3}\, I(\epsilon,y) + n(y)\, c y^{-3} \int_{\epsilon}^{b(\epsilon)} d\epsilon'\, \sigma(\epsilon|\epsilon')\, I(\epsilon',y) \tag{67}$$

where $y^{-3}\, I$ is the specific comoving photon intensity in the expanding universe, $Q_A(\epsilon, y)$ is the source term of γ-rays produced in matter-antimatter annihilations, and $n(y)$ is the average gas density in the intergalactic medium. The quantity $d\epsilon/dt$ describes the energy loss rate of photons due to the expansion of the universe which is given by

$$\frac{d\epsilon}{dt} = \frac{\epsilon}{y}\frac{dy}{dt} = -\frac{\epsilon}{y}\left[H_0\, y^2 \sqrt{1 + \Omega(y-1)}\right] \tag{68}$$

(see, e.g., McVittie 1965). The second term on the right hand side of equation (67) describes the absorption of γ-rays by intergalactic gas due to pair-production and Compton scattering. The integral term describes the transport of photons being scattered down in energy by Compton interactions.

The partial derivative

$$\frac{\partial}{\partial t} \rightarrow \frac{dy}{dt}\frac{\partial}{\partial y} = -H_0\, y^2 \sqrt{1 + \Omega(y-1)}\, \frac{\partial}{\partial y} \tag{69}$$

Substitution of equations (68) and (69) in equation (67) and subsequent simplification yields the CPT equation in the form

$$y\frac{\partial I}{\partial y} + \epsilon\frac{\partial I}{\partial \epsilon} = 2I + \frac{\Omega}{\sigma_u}\, \frac{y^2}{\sqrt{1 + \Omega(y-1)}}\Bigg[\sigma_a\, I$$

$$- \int_{\epsilon}^{b(\epsilon)} d\epsilon'\, \sigma(\epsilon|\epsilon')\, I - \xi^2 \Omega\, n_c\, G_A(\epsilon)\, y^3\, \Phi(y)\Bigg] \tag{70}$$

where

$$\sigma_u \equiv \frac{H_0}{n_c c} \simeq 10^{-23} \text{ cm}^2 \tag{71}$$

and

$$n_c = 3H_0^2/8\pi G \simeq 10^{-5} \text{ cm}^{-3}.$$

G being the Gravitational constant.

The upper limit on the Compton-scattering integral,

$$b(\epsilon) \equiv \begin{cases} \epsilon/(1-2\epsilon), & \epsilon < \frac{1}{2} \\ \\ \infty, & \epsilon \geq \frac{1}{2} \end{cases}, \tag{72}$$

and the interaction velocity function

$$\Phi(y) \equiv \frac{\sigma_A [T(y)] \; v [T(y)]}{4\pi \; \sigma_u} \tag{73}$$

and G (ϵ) is the γ-ray source term normalized per interaction as shown in figure (3).

We further define

$$\nu \equiv \frac{\pi r_0^2}{\sigma_u} = 2.5 \times 10^{-2} \tag{74}$$

where r_0 is the classical radius of the electron. We then introduce a dimensionless absorption term $A(\epsilon) = A_c(\epsilon) + A_p(\epsilon)$ to take account of both Compton scattering and pair production where

$$A_c(\epsilon) \equiv \frac{2+2\epsilon}{\epsilon^2}\left(\frac{2+2\epsilon}{1+2\epsilon} - \frac{\ln(1+2\epsilon)}{\epsilon}\right) + \frac{\ln(1+2\epsilon)}{\epsilon} - \frac{2+6\epsilon}{(1+2\epsilon)^2}, \tag{75}$$

and

$$A_p(\epsilon) \equiv \begin{cases} 3.25\times 10^{-3}\, L(\epsilon), & L(\epsilon) = (3.12\ln\epsilon - 8.07) \geq 0 \\ 0 & , L(\epsilon) < 0 \end{cases} \tag{76}$$

A dimensionless transfer function is also defined as

$$B(\epsilon|\epsilon') \equiv \frac{\epsilon^2}{(\epsilon')^4}\left[\frac{2}{\epsilon'} - \frac{2}{\epsilon} + \frac{1}{(\epsilon')^2} + \frac{1}{\epsilon^2} - \frac{2}{\epsilon\epsilon'} + \frac{\epsilon}{\epsilon'} + \frac{\epsilon'}{\epsilon}\right] \tag{77}$$

(Heitler 1954).

Using these definitions, we obtain the final form for the CPT equation as used in the numerical calculations

$$y\frac{\partial I}{\partial y} + \epsilon\frac{\partial I}{\partial \epsilon} = 2I + \frac{y^2\Omega\nu}{[1+\Omega(y-1)]^{1/2}}\left[A(\epsilon)\, I - \int_\epsilon^{b(\epsilon)} d\epsilon'\, B(\epsilon|\epsilon')\, I(\epsilon', y) - \frac{\xi^2\Omega n_c G_A(\epsilon)\, y^3\Phi(y)}{\pi r_e^2}\right] \tag{78}$$

VIII. COMPARISON OF CALCULATED γ-RAY SPECTRUM FROM COSMOLOGICAL MATTER-ANTIMATTER ANNIHILATION WITH OBSERVATIONS

The function $I_A(E_\gamma, y = 1)/\tilde{n}^{int}_{p,0}\, \tilde{n}^{int}_{\bar{p},0}$, obtained by numerical solution of the CPT equation using the method of characteristics, is shown by the solid line in figure 4. If we chose the values $\Omega = 1$ and $\xi = 3 \times 10^{-8}$, we obtain a theoretical annihilation spectrum compatible in both form and intensity to the observed

spectrum of cosmic background γ-radiation above 1 MeV as is shown in curve A of figure 5. We conclude that the observations of Vette, et al. (1970), which indicate a marked deviation from the power-law X-ray spectrum (curve X), may be evidence of the existence of antimatter on a cosmological scale. This conclusion, however, is subject to the following conditions:

(1) There must be an additional power-law background component of X-rays below 1 MeV of another origin. Our calculations of the annihilation γ-ray spectrum below 1 MeV (not presented in their entirety here) fall well below the observed power-law spectrum in that energy range

(2) The peak in the calculated annihilation spectrum near 1 MeV is caused by absorption and scattering of the γ-rays by interactions with an intergalactic medium having an average density near the critical value $n_c \simeq 10^{-5}$ cm^{-3}.

(3) There is strong evidence that at a redshift of $z \sim 2.5$ (or perhaps somewhat greater) the intergalactic medium is strongly ionized and equation (3) does not hold (Gunn and Peterson 1965, Rees 1969). At the redshift where reionization occurs there will be a rapid drop in the cross section for annihilation (see figure 1) resulting in a sharp steepening of the γ-ray spectrum at an energy $E_{\gamma,c} \sim 50$ MeV. Above $E_{\gamma,c}$ the spectrum will drop off much more rapidly than the power-law dependence $\sim E_\gamma^{-2.86}$ shown in figure 3. Assuming that the reionization of the intergalactic medium takes place at a redshift of 2.5, we obtain the dashed line shown in figure 5 for the high-energy end of the annihilation spectrum.

(4) The observational data on the γ-ray spectrum above 1 MeV do not provide enough information at present to yield a unique identification of their

origin. There have been several alternative attempts to explain the background flux as being due to the decay of neutral pions produced by cosmic-ray interactions at high redshifts (the protar hypothesis (Stecker 1969a, 1969b, 1971)), galactic electron bremsstrahlung (Rees and Silk 1970), extragalactic electron bremsstrahlung (Silk 1970) extragalactic proton bremsstrahlung (Brown 1970) and nuclear emission lines (Clayton and Silk 1969). More recent work shows that none of these alternatives except the protar hypothesis seems capable of explaining both the form and intensity of the observed flux (Vette, et al. 1970, Stecker, et al. 1971, Stecker and Morgan 1971, Jones 1971).* The form of the -ray spectrum produced as a result of cosmic-ray interactions at high redshifts is shown in the curve marked CR in figure 5. It can be seen from figure 5 that definitive measurements of the cosmic γ-ray background spectrum above 6 MeV will provide an observational test between the protar hypothesis (curve CR) and the annihilation hypothesis (curve A). The annihilation hypothesis predicts a steepening in the background spectrum above 50 MeV whereas the protar hypothesis predicts a steepening in the spectrum at about 7 GeV (Fazio and Stecker 1970). Should future investigations indicate that the background γ-radiation above 1 MeV is not due primarily to matter-antimatter annihilation,** this will place an upper limit on the quantity ξ of 3×10^{-8} or alternatively an upper limit on the product of interacting matter and antimatter densities at present of $\tilde{n}_{p,0}^{int} \tilde{n}_{\bar{p},0}^{int} \simeq 10^{-25}$ cm^{-6}.

*A suggested explanation by Sunyaev (ZhETF Pis. Red. 12, 381 (1970)) has recently come to our attention. He suggests that bremsstrahlung radiation from relativistic electrons in the nuclei of Seyfert galaxies would have a spectrum consistent with the Vette, et al. results and have a sharp cutoff above 20 MeV consistent with the upper limits at higher energies. The sharp cutoff at 20 MeV suggested by Sunyaev has been disputed by Prilutsky, Ochelov, Rozental and Shukalov (U.S.S.R. Academy of Sciences Institute for Space Research Preprint #51, 1970). The problem of

producing a high enough intensity of gamma-rays from this mechanism has been discussed by Bisnovaty-Kogan and Sunyaev (Ap. Lett. 7, 237 (1971)).

Note on recent experimental results: Trombka (private communication) has indicated that a recent reanalysis of the Ranger 3 results above 1 MeV is in agreement with the measurements of Vette, et al. Golenetsky (Moscow Seminar on Cosmic-Rays and Astrophysics 1971) has reported that new upper limits obtained by a detector aboard the low-altitude Cosmos 135 satellite for γ-rays of energies up to 3 MeV are in conflict with the measurements of Vette, et al. Daniel (12th Int'l. Conf. on Cosmic Rays, Hobart, Australia, August, 1971, OG-27) has reported on measurements between 0.25 and 4.2 MeV which appear to be consistent with the measurements of Vette, et al. Measurements made aboard Apollo 15 are now being analysed and may be available in the near future. Uncertainties in the present data indicate the need for cleaner, more sophisticated measurements in the future.

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Table 1

Experimental Cross Sections for p-$\bar{p}$ Annihilation as a Function of Incident Antiproton Kinetic Energy

$T_{\bar{p}}$, MeV	σ_A, mb	$\beta_{\bar{p}}$, cms	$\beta_{\bar{p}}, \sigma_A$, mb	Reference
25 to 40	192 ± 34	0.13	25	Loken and Derrick (1963).
45	175 ± 45	.15	26	Cork et al. (1962).
40 to 55	155 ± 27	.16	25	Loken and Derrick (1963).
55 to 80	118 ± 26	.20	24	Loken and Derrick (1963).
90	101 ± 9	.22	22	Cork et al. (1962).
145	99 ± 8	.28	28	Cork et al. (1962).
245	66 ± 6	.36	24	Cork et al. (1962).
7000	23.6 ± 3.4	~1	24	Ferbel et al. (1965).

Table 2

Average Energies and Production Rates for Various Particles Produced in p-$\bar{p}$ Annihilations at Rest

[Matsuda, 1966]

τ, particle or pair	Average energy, MeV	Calculated production rate	Experimental production rate
π	380	3.96	3.94 ± 0.33
ρ	850	2.3×10^{-1}	$(2.5 \pm 0.6) \times 10^{-1}$
ω	940	4.2×10^{-2}	$(4.5 \pm 0.7) \times 10^{-2}$
η	860	2.4×10^{-2}	$(1.4 \pm 0.5) \times 10^{-2}$
$K\bar{K}$	(K)660	3.1×10^{-2}	$(3.3 \pm 1.6) \times 10^{-2}$
$K\bar{K}^*$, $\bar{K}K^*$		20.6×10^{-3}	$(8.8 \pm 1.8) \times 10^{-3}$
$K^*\bar{K}^*$	(K*)980	3.8×10^{-3}	$(3.9 \pm 0.7) \times 10^{-3}$

Table 3

Decay Modes, Branching Ratios, and γ-Ray Multiplicities for Various Particles Produced in p-$\bar{p}$ Annihilations

Decay mode	Branching ratio, R	$\zeta_\gamma R$
$\rho^\pm \to \pi^\pm + \pi^0$	~1.00	2.0
$\omega \to \pi^+ + \pi^- + \pi^0$	.89	1.78
$\to \pi^0 + \gamma$	.10	.30
$\eta \to \gamma + \gamma$	.386	.78
$\to 3\pi^0$ or $\pi^0 + 2\gamma$	.308	1.5
$\to \pi^+ + \pi^- + \pi^0$	.250	.5
$\to \pi^+ + \pi^- + \gamma$	.055	.05
$K^\pm \to \pi^\pm + \pi^0$ [a]	.215	.43
K_0: $K_1^0 \to \pi^0 + \pi^0$	.155	.62
K_0: $K_2^0 \to \pi^0 + \pi^0 + \pi^0$	.133	.80
K_0: $K_2^0 \to \pi^+ + \pi^- + \pi^0$	.067	.1
$K^* \to K + \pi$	—	—

[a]Other γ-ray-producing decay modes have negligible branching ratios.

Table 4

Production Rates for Various Meson-Producing Channels in p-$\bar{p}$ Annihilations at Rest [From Baltay et al., 1964]

Channel	Rate, %	ζ_γ	$\zeta_\gamma R$
$\rho^0 \pi^0$	1.4 ± 0.2	2	2.8
$\rho^\pm \pi^\pm$	2.9 ± 0.4	2	5.8
$\rho^0 \pi^+ \pi^-$ [a]	$5.8 {+0.3 \atop -1.3}$	0	0
$\rho^0 \rho^0$	0.4 ± 0.3	0	0
$\rho^0 \pi^+ \pi^- \pi^0$	7.3 ± 1.7	2	14.6
$\rho^\pm \pi^\pm \pi^+ \pi^-$	6.4 ± 1.8	2	12.8
$\omega^0 \pi^+ \pi^-$ [a]	3.8 ± 0.4	2	7.6
$\eta^0 \pi^+ \pi^-$ [a]	1.2 ± 0.3	~ 2.6	3.1
$\omega^0 \rho^0$	0.7 ± 0.3	2	—
$\eta^0 \rho^0$	0.22 ± 0.17	~ 2.6	—

[a] Includes cases where $\pi^+ \pi^-$ were from ρ^0 decay.

Table 5

Exponents of Cosmological Power-Law Annihilation γ-Ray Spectra

Ω	z	$E_{\gamma 0} \simeq \frac{70 \text{ MeV}}{(1+z)}$	δ	Γ
1	0–150	500 keV–70 MeV	0.64	2.86
	150–1000	70 keV–500 keV	0.64	2.68
	10^3–10^6	< 70 keV	2	2.00
$\simeq 0$	0–150	500 keV–70 MeV	0.64	3.36
	150–1000	70 keV–500 keV	0.64	3.18
	10^3–10^6	< 70 keV	2	2.50

FIGURE CAPTIONS

Figure 1. Proton-antiproton annihilation cross sections as a function of temperature and kinetic energy for free nucleons and atoms (Morgan and Hughes, 1970). The curve marked σ_p takes into account the mutual coulomb attraction between the proton and antiproton; the curve marked σ_p is an extrapolation from the accelerator data that fails to take this effect into account.

Figure 2. Cross section times velocity for hydrogen-antihydrogen annihilation of nucleons given as a sum for the interactions shown in figure 1 integrated over a Maxwell-Boltzmann distribution at temperature T and taking ionization into account. (a) $\sigma_{H+\bar{H}}$ (extrapolated). (b) $\sigma_{H+\bar{H}}$ (calculated). (c) $\sigma_{H+\bar{H}}$, $\sigma_{p+\bar{H}}$, $\sigma_{\bar{p}+H}$ and $\sigma_{p+\bar{p}}$, multiplied by appropriate factors to take account of fractional ionization as determined using the Saha equation. (d) $\sigma_{p+\bar{p}}$, taking coulomb attraction into account. (e) $\sigma_{p+\bar{p}}$, from high-energy accelerator data.

Figure 3. Normalized γ-ray spectrum from p-$\bar{p}$ annihilation.

Figure 4. The cosmological γ-ray annihilation spectrum calculated by numerical solution of the CPT equation (equation (80)) for $\Omega = 1$. The solid line represents the complete solution. The other curves show the effect of neglecting the absorption and scattering terms in the CPT equation.

Figure 5. Observations of the γ-ray background spectrum together with the theoretical spectra discussed in the text.

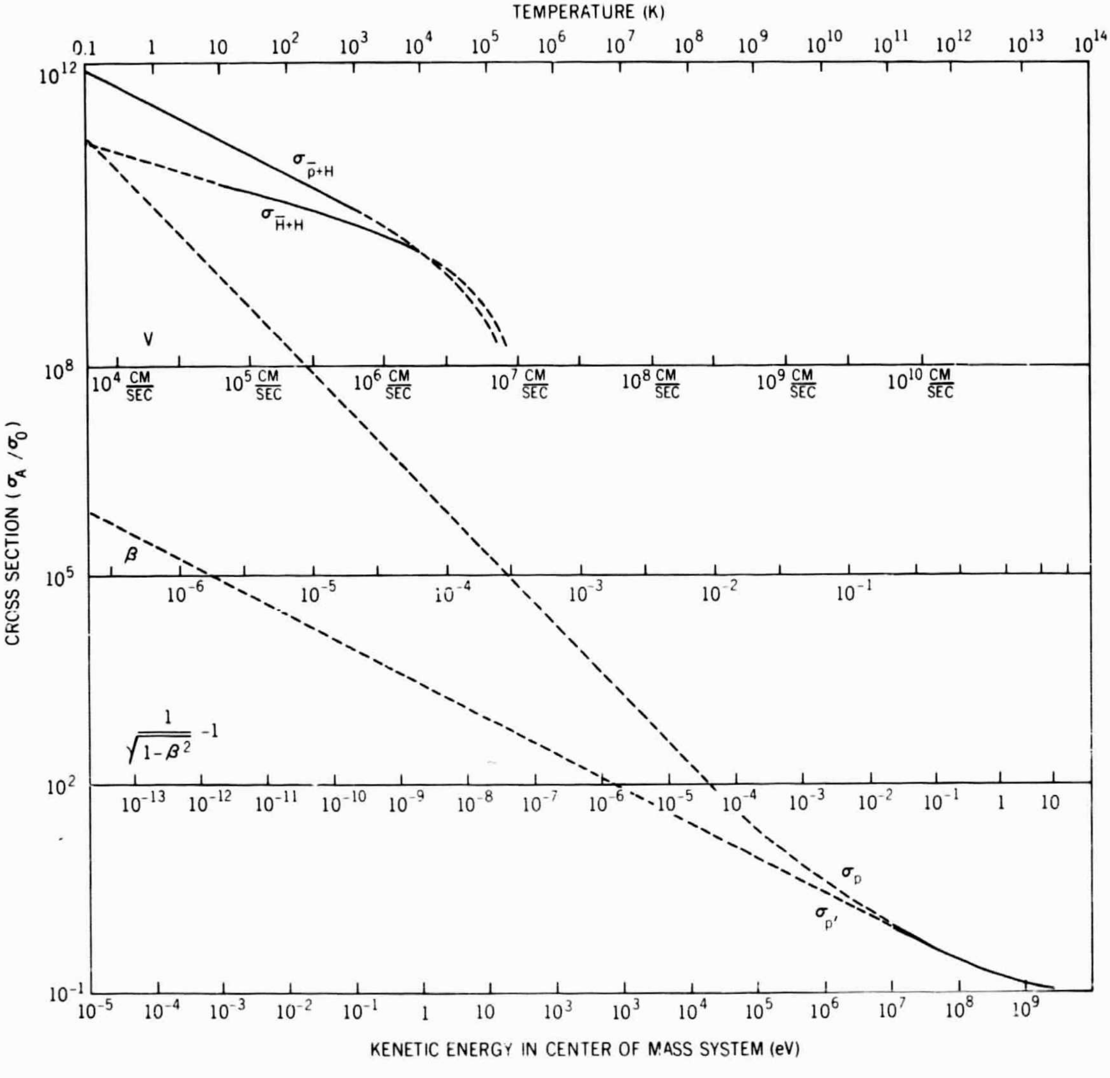


Figure 1. Proton-antiproton annihilation cross sections as a function of temperature and kinetic energy for free nucleons and atoms (Morgan and Hughes, 1970). The curve marked σ_p takes into account the mutual coulomb attraction between the proton and antiproton; the curve marked $\sigma_{p'}$ is an extrapolation from the accelerator data that fails to take this effect into account.

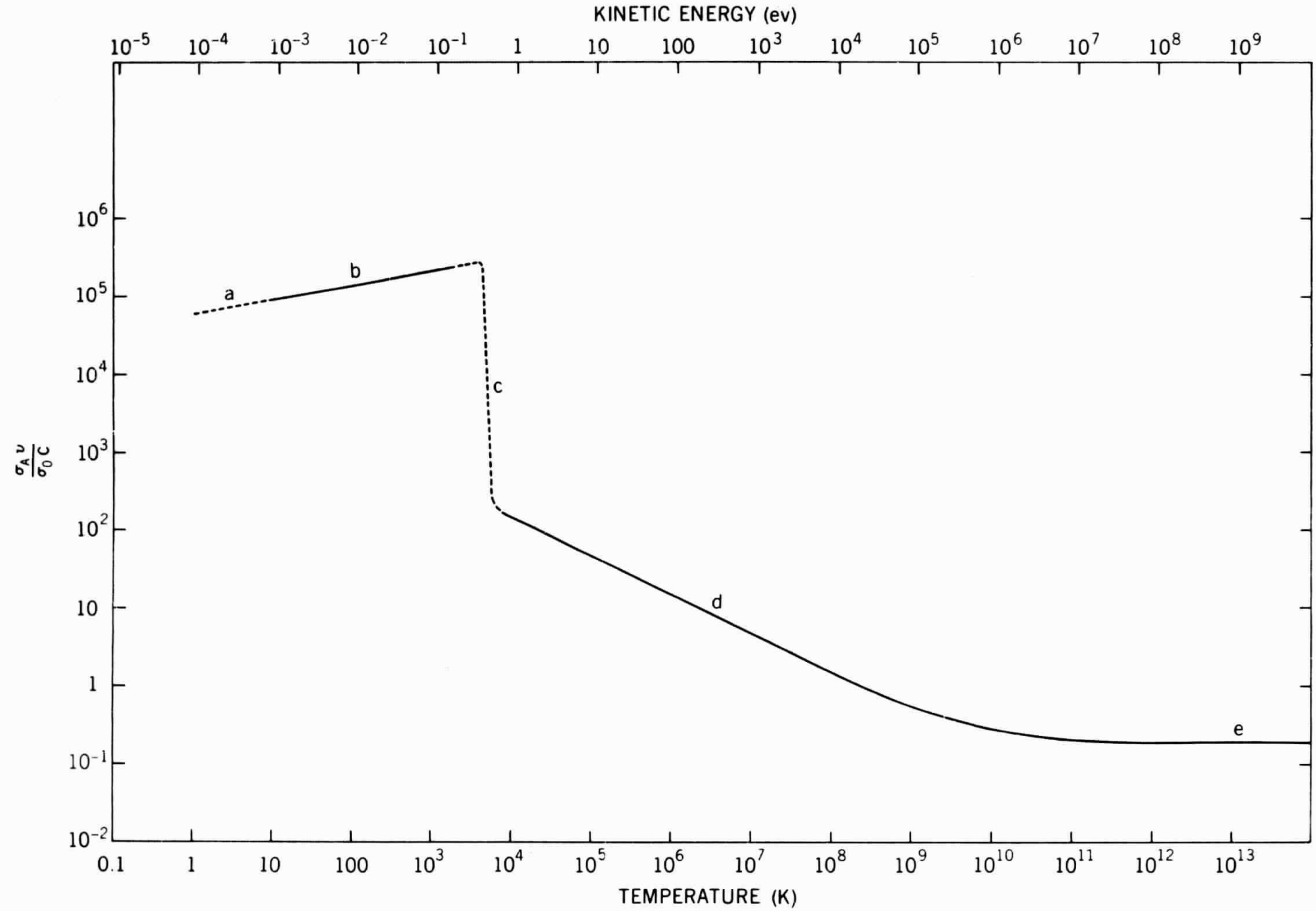


Figure 2. Cross section times velocity for hydrogen-antihydrogen annihilation of nucleons given as a sum for the interactions shown in figure 3-1 integrated over a Maxwell-Boltzmann distribution at temperature T and taking ionization into account. (a) $\sigma_{H+\bar{H}}$ (extrapolated). (b) $\sigma_{H+\bar{H}}$ (calculated). (c) $\sigma_{H+\bar{H}}$, $\sigma_{p+\bar{H}}$, $\sigma_{\bar{p}+H}$ and $\sigma_{p+\bar{p}}$, multiplied by appropriate factors to take account of fractional ionization as determined using the Saha equation. (d) $\sigma_{p+\bar{p}}$, taking coulomb attraction into account. (e) $\sigma_{p+\bar{p}}$, from high-energy accelerator data.

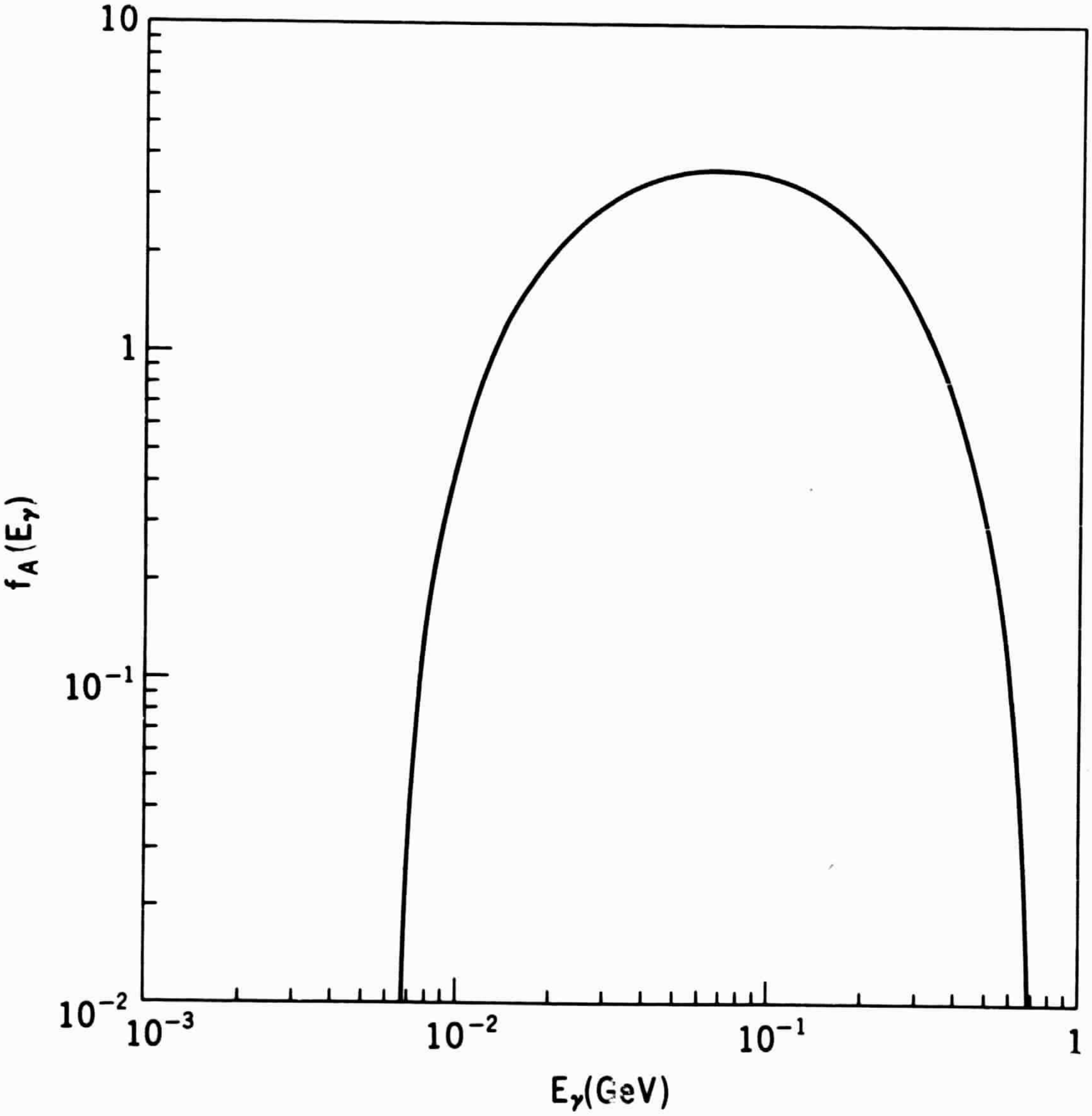


Figure 3. Normalized γ-ray spectrum from p–$\bar{p}$ annihilation

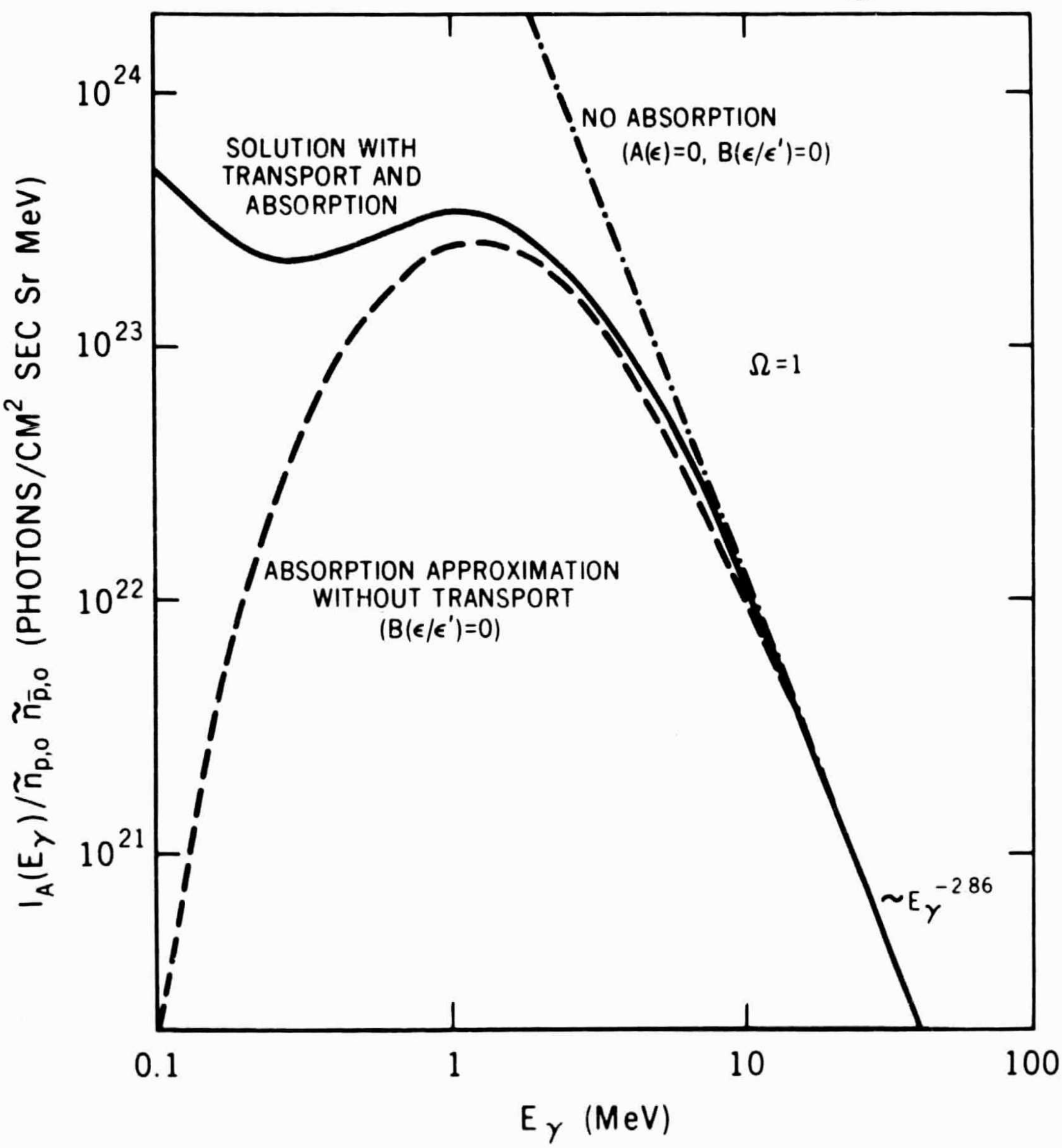
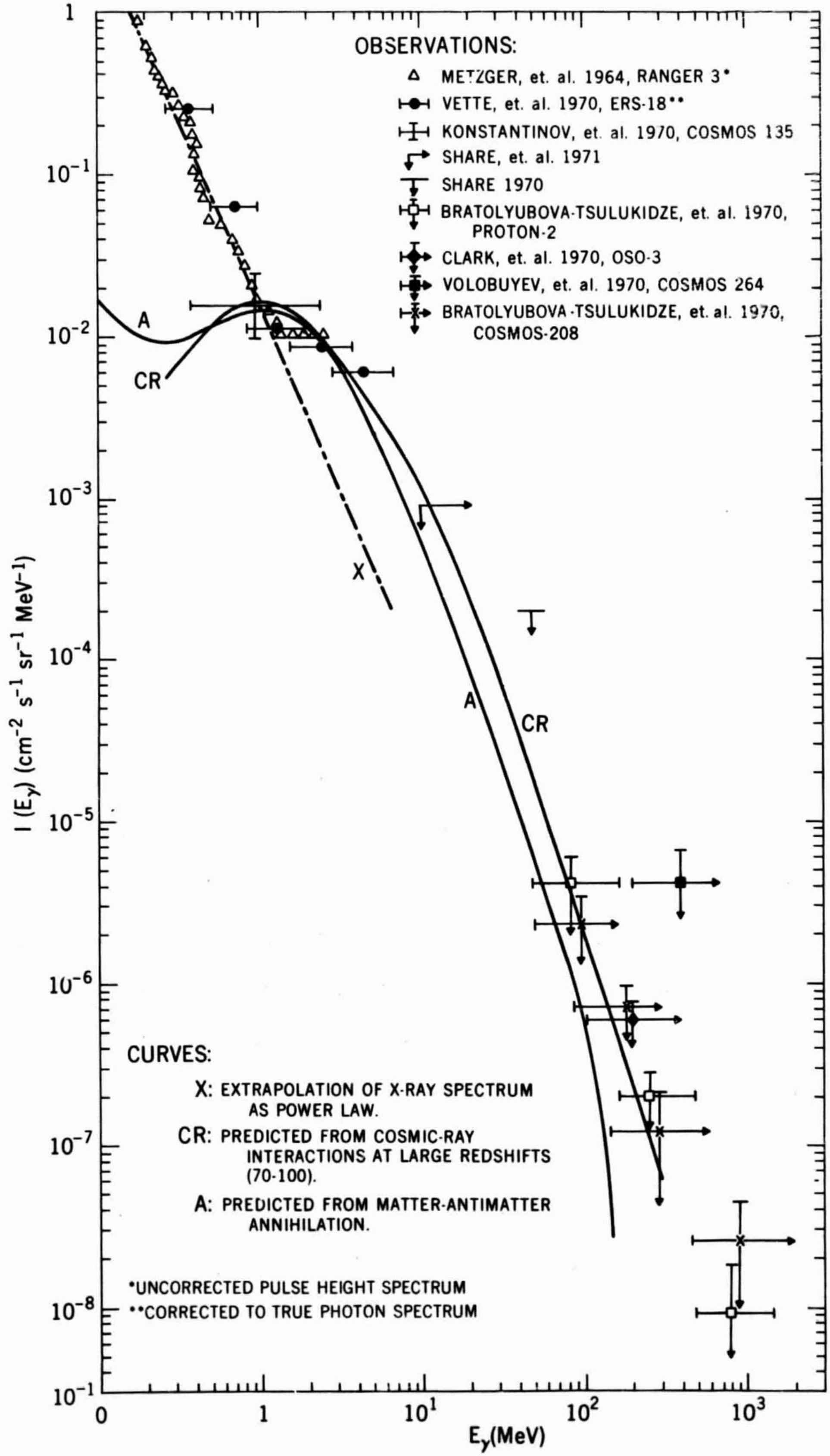


Figure 4. The cosmological γ-ray annihilation spectrum calculated by numerical solution of the CPT equation (equation (80)) for $\Omega = 1$. The solid line represents the complete solution. The other curves show the effect of neglecting the absorption and scattering terms in the CPT equation.


OBSERVATIONS:
METZGER, et. al. 1964, RANGER 3*
VETTE, et. al. 1970, ERS-18**
KONSTANTINOV, et. al. 1970, COSMOS 135
SHARE, et. al. 1971
SHARE 1970
BRATOLYUBOVA-TSULUKIDZE, et. al. 1970, PROTON-2
CLARK, et. al. 1970, OSO-3
VOLOBUYEV, et. al. 1970, COSMOS 264
BRATOLYUBOVA-TSULUKIDZE, et. al. 1970, COSMOS-208
A
CR
X
A
CR
CURVES:
X: EXTRAPOLATION OF X-RAY SPECTRUM AS POWER LAW.
CR: PREDICTED FROM COSMIC-RAY INTERACTIONS AT LARGE REDSHIFTS (70-100).
A: PREDICTED FROM MATTER-ANTIMATTER ANNIHILATION.
*UNCORRECTED PULSE HEIGHT SPECTRUM
**CORRECTED TO TRUE PHOTON SPECTRUM
$I(E_\gamma)$ (cm^{-2} s^{-1} sr^{-1} MeV^{-1})
E_γ(MeV)
1
10^{-1}
10^{-2}
10^{-3}
10^{-4}
10^{-5}
10^{-6}
10^{-7}
10^{-8}
10^{-1}
0
1
10^1
10^2
10^3

APPENDIX: THE BARYON SYMMETRIC COSMOLOGY OF OMNES AND THE VALUE OF ξ

J. L. Puget
F. W. Stecker

1. The Omnes Model

According to the calculations of Omnes, (1971a,b), in black body radiation at high temperature (≥ 350 Mev) there is a phase transition which leads to the separation of matter and antimatter. In the big bang model of the universe this corresponds to a time 10^{-5} sec after the initial singularity.

During the period when the temperature remains higher than a fraction of an Mev (until the end of the leptonic era) the nucleons and antinucleons diffuse. The annihilation is only limited by the speed of diffusion and the statistical fluctuations which depend upon the size of the initial separation. The next period is called by Omnes the coalescence period. The study of the behavior of an emulsion of matter and antimatter by Omnes (1971a,b) shows that the characteristic size of a region of pure matter (or antimatter) in the emulsion is independent of the initial mechanism and only depends on the ratio of the density of nucleons to the density of black body photons as a parameter. This ratio $\eta = n_p/n_{ph}$ is 10^{-8} for $\Omega = 1$.

The coalescence period ends at a temperature $\sim\frac{1}{4}$ eV when the plasma combines, and the size of the emulsion regions is fixed by the time at which the plasma forms neutral atoms and is $\sim\lambda_\gamma$ where λ_γ is the mean free path of the γ rays at that time.

2. The Neutral Period

Omnes (1971a) considers that the start of the neutral period in his cosmology occurs at the same time as in the pure matter big bang model. We will now show that this is very inaccurate because the annihilation process results in ionizations which delay the start of the neutral period. Peebles (1968) has given a detailed solution of the combination problem when the plasma is only composed of matter. The equation he gives for the electron density is:

$$\frac{d}{dt}\, n_e = \left(\alpha_c\, n_e^2 - \beta_c\, n_H\, e^{-\frac{B_1 - B_2}{kT}} \right) C(T)$$

where B_n is the binding energy for hydrogen with the electron in an orbit corresponding to the n^{th} principal quantum number and n_H is the density of hydrogen atoms in the ground state. The α_c, β_c and C (T) coefficients are functions of the temperature T.

The annihilation rate per unit volume is $\rho_a = \frac{\nu_a}{R}$ where R is the characteristic dimension and ν_a is the annihilation rate per unit surface area of the annihilation layer,

$$\nu_a = h\, n_p\, n_{\bar{p}} \langle \sigma v \rangle$$

with h given by

$$h = \left[24 \frac{D}{N \langle \sigma v \rangle} \right]^{1/2}$$

and D is the diffusion coefficient.

This diffusion coefficient has been discussed by Omnes (1971a). He finds

$$\nu_a = N\, 0.48\, v_t$$

where v_t is the thermal velocity and N is the density of nucleons (antinucleons) at the boundary of the annihilation layer. N is lower than n_p (or $n_{\bar{p}}$) because of the effect of the annihilation pressure (Omnes 1971b). During the coalescence period the ratio $\frac{N}{n_p}$ is almost constant and is $\simeq 0.05$.

In order to investigate the effect of this annihilation on the recombination of the plasma one must consider the behavior of all the products of the annihilation. The neutrinos escape carrying $\simeq 50\%$ of the energy. The γ-rays lose their energy by Compton collisions with the electrons of the medium with an effective cross section $\langle\sigma_\gamma\rangle \simeq \frac{1}{60}\sigma_T$ (σ_T is the Thompson cross section). The secondary electrons have an energy of the same order of magnitude as the γ-rays. γ-rays are also produced by annihilation of the electrons and positrons after they have been slowed down.

The electrons and positrons undergo two different types of interactions; inverse Compton interactions on the black body radiation with a cross section σ_T, and collisions with other electrons (free or bound). For a value of the parameter $\eta = 10^{-8}$ and for $z \simeq 1000$ the first process is dominant for relativistic electrons. Using a simple representation of the spectrum of the electrons produced in the annihilation: 0.86 electron with an energy 40 MeV, 0.41 with an energy 135 MeV and 0.1 with an energy 300 MeV, one gets a rough idea of the number and the energy of the X-rays produced for one annihilation: 10^4 with the energy 20 keV, 2.10^4 with the energy 5 keV and 5.10^4 with the energy 0.5 keV. This estimate is based on the assumption that the average energy of the X-rays produced in a black body radiation at the temperature T by relativistic electrons of energy E is

$$\epsilon_x = k\,T \left(\frac{E}{m\,c^2}\right)^2 \quad \text{(e.g. Ginzburg 1969)}$$

These X-rays can themselves undergo two interaction processes, photoionization and Compton collisions with free electrons. The photoionization cross section depends strongly upon ϵ_x

$$\sigma_x = \sigma_T \left(\frac{2\sqrt{2}}{137^4}\right)\left(\frac{m\,c^2}{\epsilon_x}\right)^{7/2}$$

(see, e.g., Heitler 1954)

These X-rays are very efficient ionizing agents and tend to slow down the recombination of the plasma. At the beginning of their life, the most energetic X-rays are degraded by Compton collisions with the free electrons. This process becomes unimportant when the energy of the X-rays falls below a few keV at which point photoionization becomes important if only one percent of the hydrogen is neutral. In this way, half of the energy released in the annihilation, which is carried by the γ-rays and the electrons goes into thermal electrons in a short time compared to the cosmological time.

The number of ionizations per annihilation is thus given by

$$k_i = \frac{m_p\, c^2\,(\mathrm{eV})}{32}\;\frac{n_H}{n_p}$$

because one ionization corresponds to a loss of 32 eV by a fast electron.

The equation for the recombination equilibrium then becomes

$$\frac{d}{dt}\, n_e = \alpha_c\, n_e^2\, C(T) - \rho_a \cdot k_i = 0$$

which reduces to

$$C(T, x)\, x^2 \;=\; 2.5\; 10^{-2}\, \frac{T}{10^4}\, (1 - x)$$

where

$$x \;=\; \frac{n_e}{n}\;.$$

It follows from Table 1 that the recombination is very slow compared to the big-bang model without matter-antimatter symmetry. It follows that the final size of the matter and antimatter regions is correspondingly larger than that calculated by Omnes (1971a).

However, two other effects must be considered in determining the final size of these regions:

1. E. Schatzman (1970) has shown that even very small magnetic fields can drastically alter the annihilation rate. When matter and antimatter annihilate on an interface a natural mechanism leads to the building of a magnetic field (see below).

2. If the plasma remains ionized much longer than Omnes (1971a) supposed, it is not certain that the annihilation provides enough momentum for the coalescence to continue until neutralization is complete.

We will not deal here with these problems which we hope to discuss later. Omnes (1971a) concluded that the final mass of these regions would be of the order of large galaxies. We conclude that it must be larger because of the above discussion and can be of the order of galaxy clusters.

This conclusion is supported by the fact that the collisions of galaxies and antigalaxies in dense clusters, if we suppose their number equal, leads to a γ-ray flux above 100 MeV larger than the Clark et al. (1968) limit by a factor 10^3, (Puget 1971, unpublished).

3. Metagalactic Annihilation As a Function of the Red-Shift

3.1 Initial Conditions

At z somewhere between 300 and 1000 the effect of neutralization and of annihilation pressure is to stop the coalescence by separating matter and antimatter into independent clouds. We suppose that these clouds are the protoclusters of galaxies. We know their actual number, $\mathfrak{N}_0$, is $\sim 5 \times 10^{-5}$ Mpc^{-3} (Abell 1965). Denoting the red-shift corresponding to the time of separation of the clouds as a parameter z_i, we have computed the radius of the clouds at this time, and their corresponding mass as determined by the relation

$$\frac{4}{3}\pi R^3 \mathfrak{N}_0 (1 + z_i)^3 = 1$$

where $R(z_i) \simeq \frac{1}{2} \lambda_\gamma (z_i)$ is the final radial dimension of the separate regions as previously discussed. The results are shown in Table 2.

3.2 Cloud-Anticloud Collisions

For $z < 200$ the clouds are well separated and the annihilation rate is directly related to the collision between clouds and anticlouds. When a cloud of neutral matter collides with a cloud of neutral antimatter, all the atoms from one cloud reaching the other cloud are annihilated; the rate of annihilation is given by $\nu_a = n_H v$ (v is the relative speed). The strong flux of electrons and γ-rays produces by Compton interactions a current of the same sign on both

sides of the annihilation layer and one can see that, for the collision of two roughly spherical clouds, this current produces a very low primary magnetic field but one which is strong enough so that the Larmor frequency is greater than the inverse electron lifetime by Compton losses. (Puget 1971, unpublished) Under these conditions, the current is not normal to the annihilation layer but parallel to it and we have to consider the behaviour of the magnetic field in relation to this drift current as the ionization front propagates. The idea of the following computations is due to E. Schatzman (1970). The general process remains the same but some details have been corrected.

Figure 1 depicts the processes involved. Two ionization fronts are moving symmetrically with speed c on both sides of the annihilation layer.

This ionization is mainly due to photoionization by the X-rays. The electrons ejected in this process have an initial energy $\sim\epsilon_x$ and produce a drift current when they are slowed down. All the characteristic times for these processes are shorter than τ, the lifetime of the primary electrons produced in the annihilation. (τ is also the characteristic time for the production of the X-rays.)

Between the annihilation layer and the ionization front, a flux of X-rays and a flux of γ-rays produce fast electrons by Compton collisions. These fast electrons, when slowed down, generate a drift current $\vec{J}$ of opposite sign on each side of the annihilation layer.

In the plasma, the particles are in equilibrium between the force per unit volume F_x due to the X- and γ-rays and the electromagnetic force per unit volume

$$F_x = -\frac{\vec{J} \times \vec{B}}{c} - \rho \vec{E}$$

Maxwell's equations together with the fact that all the field quantities are independent of y due to the choice of the frame of reference (indicated in Fig. 1) yield the relation

$$\frac{\partial Bz}{\partial x} + \frac{v_X \partial Bz}{c^2 \partial t} = 4\pi \frac{F_x}{B_z} \qquad \text{(A-1)}$$

if we suppose that the drift velocity of the plasma, v_X, is almost independent of t.

The force F_x has two components coming from the X-rays and the γ-rays.

$$F_x = \sum_{\epsilon_\gamma} \nu_a \sigma_\gamma \frac{n_e}{c} n_\gamma \epsilon_\gamma + \sum_{\epsilon_x} \nu_a \sigma_\gamma \frac{n_e}{c} n_x \epsilon_x$$

where n_γ is the number of γ-rays of energy ϵ_γ produced in one annihilation. And where we have used a simple representation of the annihilation γ-ray spectrum assuming

50% production with an energy of 70 MeV

35% production with an energy of 250 MeV

and 15% production with an energy of 500 MeV,

to estimate the value of F_x.

The solution of the Equation (A-1) gives

$$\frac{B^2}{8\pi} \approx cF_X(x + v_X t) \quad \text{for } x < ct - c\tau$$

and

$$\frac{B^2}{8\pi} \approx \left(\frac{B^2}{8\pi}\right)_{(ct-c\tau)} + F_x' v_X \tau \, [1 - \kappa^2] \quad \text{for } ct - c\tau < x < ct$$

where

$$\kappa = \frac{-x + ct}{c\tau}$$

and F_x' is the force in the ionization front, mainly due to photoionization process of the hydrogen by the X-rays,

$$F_x' = \int \phi \, n_H \, \sigma_x \, \frac{2\,\epsilon_x}{c} \, d\,\epsilon_x$$

where ϕ is the differential flux of X-rays.

The quantity $\frac{B_z^2}{8\pi}$ as a function of x is given on Figure 2.

The time necessary for developing a magnetic field strong enough so that

$$\frac{B^2}{8\pi} \geq \frac{1}{2} n \, m_p v^2$$

is given by

$$t_B = \frac{1}{2} n \, m_p v^2 \, \frac{1}{F_x \, c}$$

This time is much larger than τ and justifies the approximations made. The field and the charged particles move together. When the energy density in the field in one cloud is much larger than the kinetic energy of the other cloud, the field cannot propagate to the other side of the annihilation layer. The annihilation stops and the collision becomes almost elastic with a "magnetic cushion" between the two clouds. The surface area of the two clouds in contact is a function of the impact parameter, b, and the radius, R given by

$$s \simeq \frac{\pi(2R - b)}{2}\left(R^2 - \frac{b^2}{4}\right)^{1/2} \qquad \text{(A-2)}$$

where R is the radius of the clouds. The annihilated mass is

$$M_a = \frac{1}{2} n^2 m_p^2 v^3 s \frac{1}{F_x c}.$$

The fraction

$$\zeta \equiv \frac{M_a}{M_t} = \frac{1}{M_t} \frac{m_p^2 v^2}{8 \times 10^{-28}} s \qquad \text{(A-3)}$$

The number of these clouds per unit volume is $\mathfrak{N} = \mathfrak{N}_0 (1 + z)^3$. $\mathfrak{N}_0$ is their actual density can be evaluated as: $5\ 10^{-5} > \mathfrak{N} > 10^{-5}$ (Mpc^{-3}) depending on the value of the Hubble constant.

The average density of matter in the universe is given by

$$n(z) = n_0 (1 + z)^3$$

so the mass of the clouds is fixed by the value of $\mathfrak{N}_0$ and the hypothesis $\Omega = 1$. The total mass $M_t = 2.8\ 10^{48}$ g and $1.4\ 10^{49}$ g with $\mathfrak{N}_0$ given by 5×10^{-5} and 10^{-5} Mpc^{-3} respectively. The relative cloud velocity v is chosen to be 3×10^8 cm/s (Abell, 1961). During most of the collision period, R is almost constant because of the strong cooling by the black body radiation. R is certainly increasing in a non-negligible manner for low z (< 10) but other changes appear at the same time: strong decrease of H - $\bar{H}$ cross-section (Bardsley 1971), formation of galaxies, existence of a magnetic field created in the previous collisions and very low probability for intercloud collision. We conclude that

the behaviour of the annihilation rate at low red-shift is not clear and the formula given below can be wrong by a large factor at very low red-shift.

3.3 Annihilation Rate as a Function of z.

The number of collisions is given by:

$$dN = \frac{1}{4} \mathfrak{N}^2 \, 2\pi b db \, v$$

where b is the impact parameter.

The annihilation rate is:

$$\rho_a = \int_0^{2R} \frac{\mathfrak{N}^2}{4} \, 2\pi b db \, v \, \frac{M_a}{m_p} \tag{A-4}$$

where $M_a = M_a \, [\zeta(b)]$.

Using Equation (A-3), we can define

$$<\zeta> = \int_0^{2R} db \, \frac{2\,b}{4\,R^2} \, \zeta(b)$$

and rewrite Equation (A-4) in the form

$$\rho_a = \pi \, \mathfrak{N}_0^2 \, R^2 \, v <\zeta> \frac{M_t}{m_p} (1 + z)^6$$

or

$$\rho_a = \frac{n_{p,0}}{\tau_{c,0}} <\zeta> (1 + z)^6 \tag{A-5}$$

where $\tau_{c,0} \equiv [\mathfrak{N}_0 \pi \, R^2 \, v]^{-1}$ is the collision time for the clouds.

4. Conclusion and Outlook

We now compare these results with the assumptions of the main paper. From Equation (A-5) it follows that the annihilation rate is proportional to $(1 + z)^6$. As we have previously shown, R is almost constant for $z > 10$. Furthermore any possible increase of R during the expansion would imply a decrease in v. Considering that these quantities cannot change by more than a factor 10 we note that their product does not decrease more with decreasing z than the factor $\langle\sigma v\rangle$ of the main paper.* The value of the parameter ξ taken as 3×10^{-8} in the main paper was chosen to fit observational data. The annihilation rate used in the main paper is:

$$\rho'_a = \tilde{n}_{p,0}\, \tilde{n}_{\bar{p},0} \langle\sigma v\rangle (1 + z)^6$$

which can be written in the same form as Equation (A-5) as

$$\rho'_a = n_{p,0}\, \xi^2 \frac{1}{\tau_{a,0}} (1 + z)^6 \qquad \text{(A-6)}$$

where $\tau_{a,0} \equiv (n_{\bar{p},0}\, \sigma v)^{-1}$ is the annihilation time. So that the relation between ξ and ζ is obtained from Equations (A-5) and (A-6) as

$$\frac{\xi^2}{\tau_{a,0}} = \frac{\zeta}{\tau_{c,0}}$$

with $\tau_{a,0} \simeq 10^{14}$ sec and thus

$$\xi \simeq 10^7 \sqrt{\frac{\zeta}{\tau_{c,0}}} \qquad \text{(A-7)}$$

Abell (1961, 1965) and Schatzman (1970) indicate that v is between 10^8 and 5.10^8 cm sec^{-1} and we know that v is decreasing with time. The possible initial

*In the main paper it was assumed that the interaction velocity is the thermal velocity of the gas. This assumption results in a slow dependence of $\langle\sigma v\rangle \sim \phi(1+z) \sim (1+z)^{0.36}$. Although the model considered here is physically more plausible, it results in a similar slow dependence of $R^2 v_\zeta$ on z such that $R^2 v_\zeta \sim (1+z)^\alpha$ where $0<\alpha<1$ so that both assumptions lead to similar results as given in figure 4 to within the uncertainties of the cosmological parameters used.

values of R are given in Table 1 and those which fulfill the relation $R \leq \frac{1}{2} \lambda_\gamma$ are in the range 10^{23} cm to 5 10^{23} cm. Evaluation of Equation (A-3) then indicates that $10^{-8} \leq \zeta \leq 10^{-6}$. The average values of $\tau_{c,0}$ are 1.5 10^{22} sec and 6.7 10^{22} sec respectively for $\mathfrak{N}_0 = 10^{-5}$ Mpc^{-3} and 5 10^{-5} Mpc^{-3}. Using (A-7) we obtain the corresponding values of

$$\xi \simeq \begin{cases} 3 \times 10^{-8} & \text{for } \mathfrak{N}_0 = 10^{-5}\ \text{Mpc}^{-3} \\ \\ 1 \times 10^{-8} & \text{for } \mathfrak{N}_0 = 5 \times 10^{-5}\ \text{Mpc}^{-3} \end{cases}$$

These two values are known with an uncertainty of a factor 5, thus both are consistent with the value given in the main paper. We thus conclude that the baryon symmetric cosmology of Omnes can lead to a cosmological γ-ray spectrum of the same form and intensity as that observed.

*G. Steigman (1969) has claimed that the γ-rays experiments rule out the possibility of equal amounts of matter and antimatter in the universe. We have already shown that this is not the case.

More recently Jones and Jones (1970) arrived at the same conclusion as Steigman by considering the problem of the collisions. Their main assumption is that half of the mass is annihilated (after forgetting the existence of electrons in the annihilation products, and neglecting without justification the magnetic effects). It follows from Equations (A-2) and (A-3) that the average ratio of the mass annihilated per collision to total mass can be estimated by

$$< \zeta > \simeq \frac{1}{M_t} \frac{m^2 v^2 R^2}{2 \times 10^{-24}}$$

The order of magnitude of $< \zeta >$ is $\sim 10^{-7}$.

Sunyaev and Zeldovich (1970) have considered the question of existence of antimatter on large scale in the universe and have shown that, for $\Omega = 1$, less than 10^{-6} of the matter must have been annihilated during the period $0 < z < 100$. This condition is fulfilled in the model discussed here.

A model considered previously by one of us (Puget 1971) where diffusion plays the main role in cosmological annihilation. The results presented here indicate that cloud collisions dominate over diffusion in the redshift range of significance for γ-ray production so that we now consider the collision model more physically plausible than the diffusion model.

APPENDIX

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Table 1: Neutralization of the Cosmic Plasma

T	Peebles (1968)		This Work	
	x	$C_{(T,x)}$	x	$C_{(T,x)}$
4500	0.92	$1.8\ 10^{-4}$	0.96	$3.6\ 10^{-4}$
4000	0.40	$5.9\ 10^{-4}$	0.84	$2.2\ 10^{-3}$
3500	0.072	$2.7\ 10^{-3}$	0.64	$7\ 10^{-3}$
3000	$9.8\ 10^{-3}$	$2\ 10^{-2}$	0.38	$3.2\ 10^{-2}$
2500	$9.2\ 10^{-4}$	0.25	0.14	0.3
2000	$1.23\ 10^{-4}$	0.96	0.068	1
1500	$5.3\ 10^{-5}$	1	0.06	1
1000			0.049	

Table 2: Size of the Clouds When They Separate

z_i	n (cm^{-3})	$\mathcal{N}$ (Mpc^{-3}) $(\mathcal{N}_0 = 5\ 10^{-5})$	R (cm)	λ_γ (cm)	M (g)
1000	$3\ 10^{+3}$	$5\ 10^{4}$	$5.1\ 10^{22}$	$3\ 10^{22}$	$2.8\ 10^{48}$
600	$7\ 10^{2}$	$1.08\ 10^{4}$	$8.5\ 10^{22}$	$1.4\ 10^{23}$	$2.8\ 10^{48}$
300	10^{2}	$1.35\ 10^{3}$	$1.7\ 10^{23}$	10^{24}	$2.8\ 10^{48}$
150	10	$1.7\ 10^{2}$	$3.3\ 10^{23}$	10^{25}	$2.8\ 10^{48}$
		$(\mathcal{N}_0 = 10^{-5})$			
1000	$3\ 10^{3}$	10^{4}	$8.6\ 10^{22}$	$3\ 10^{22}$	$1.4\ 10^{49}$
600	$7\ 10^{2}$	$2.16\ 10^{3}$	$1.5\ 10^{23}$	$1.4\ 10^{23}$	$1.4\ 10^{49}$
300	10^{2}	$2.7\ 10^{2}$	$2.8\ 10^{23}$	10^{24}	$1.4\ 10^{49}$
150	10	$3.4\ 10^{1}$	$5.6\ 10^{23}$	10^{25}	$1.4\ 10^{49}$

FIGURE CAPTIONS

Figure 1. Geometry of Intercloud Collision: The interface between the two clouds is the y, z plane. The relative motion of the clouds is along the x axis. The **B** field points downward in the z direction and the effective current **J** is in the y direction. The electrons emitted in the annihilation and those accelerated by Compton collisions with γ-rays loose their energy by collision with the photons of the black body radiation generating X-rays.

Figure 2. The energy-density of the magnetic field at a time t after the beginning of the collision in either of the two clouds, as a function of x.

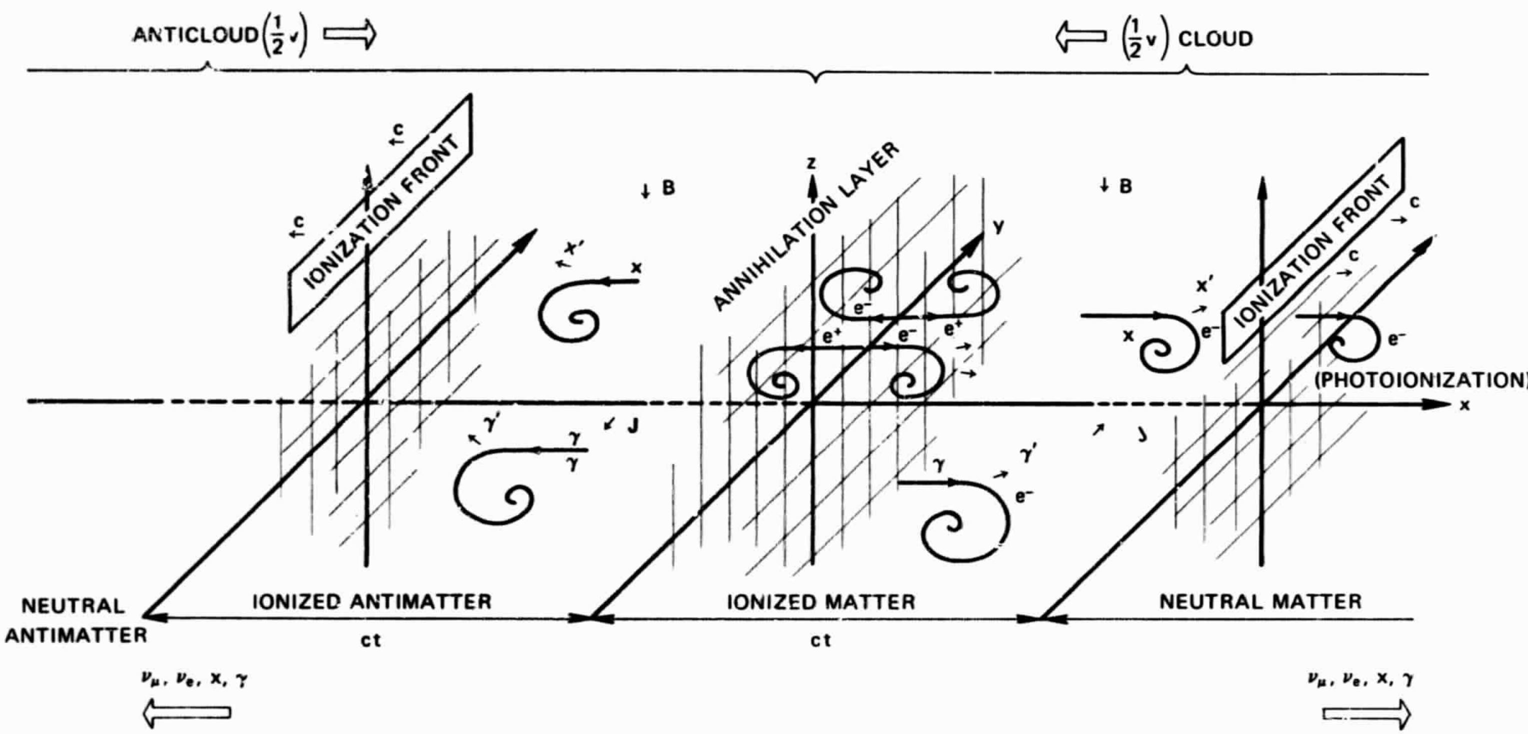

Figure 1. Cloud-Anticloud Collision, Generation of the Magnetic Field

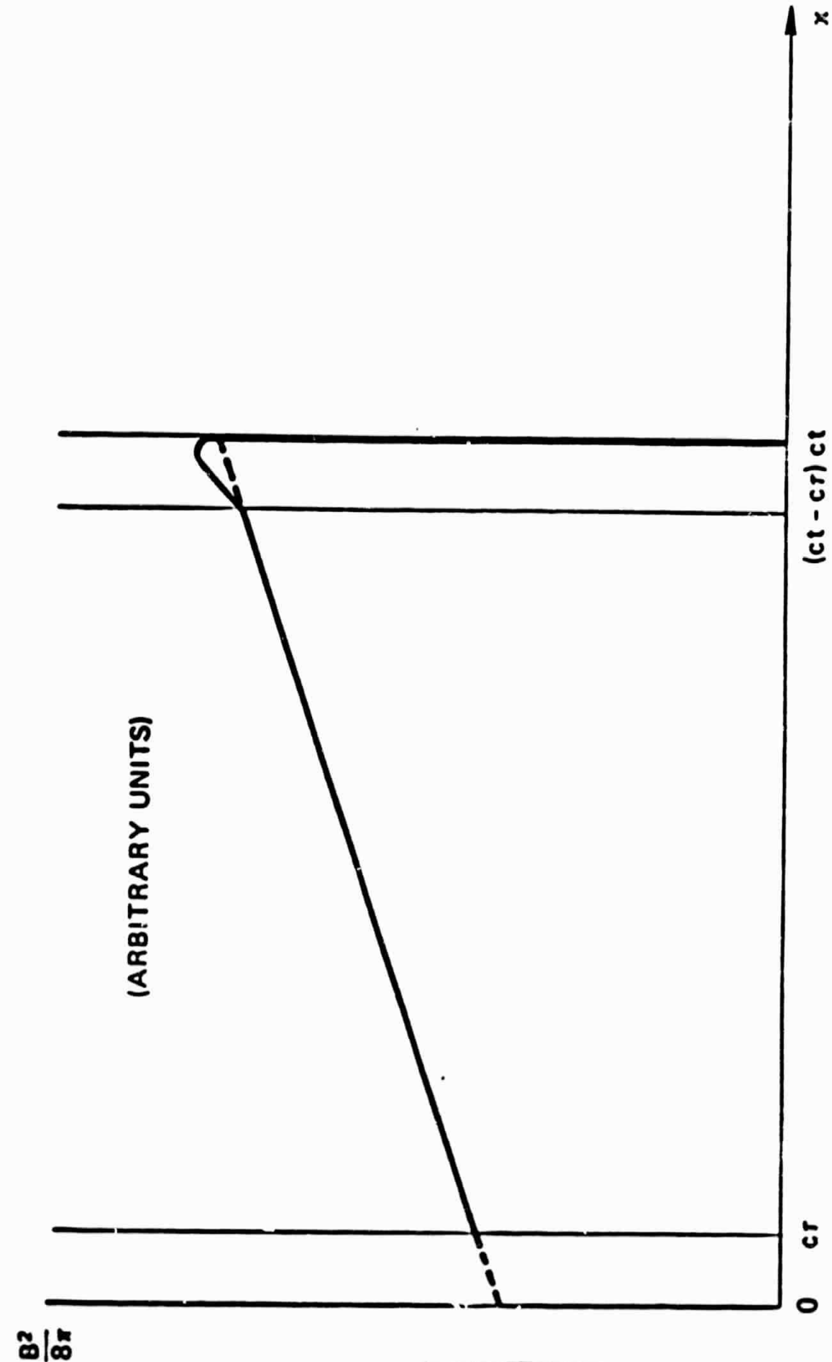

Figure A-2.